

Constraints on Inflation Models from Supersymmetry Breakings

Koshiro Suzuki and T. Watari

*Department of Physics, University of Tokyo,
Tokyo 113-0033, Japan
(February 1, 2008)*

Abstract

Effects of soft SUSY-breaking terms on inflation potentials are discussed. There exist generic constraints that must be satisfied in order not for the inflaton potential to lose its flatness. We examine explicitly the constraints in the case of a hybrid inflation model and find that the coupling constant λ between the inflaton and the “water-fall-direction” field is bounded as $2.0 \times 10^{-6} \lesssim \lambda$. This is a highly non-trivial result. Indeed, if we adopt the severest constraint from avoiding the problem of the gravitinos produced non-thermally, $\lambda \lesssim 7.4 \times 10^{-6}$ is required, under a reasonable assumption on the reheating process. This means that the hybrid inflation model marginally has a viable parameter space. We also discuss analogous constraints on other inflation models.

Supersymmetry (SUSY) is not only motivated from the low-energy particle physics, but also from the primordial inflation. SUSY naturally explains the existence of the inflaton scalar field, and moreover, the cancellation of radiative corrections in supersymmetric theories is preferable for maintaining the flatness of the inflaton potential. There have been constructed various models for inflation in the framework of supergravity such as chaotic [1], hybrid [2], new [3], topological [4], D-term [5] inflation models and so on.

Each model, in general, has some free parameters even after the COBE-normalization condition

$$\frac{\delta\rho}{\rho} = \frac{1}{5\sqrt{3}\pi} \frac{V^{3/2}(\varphi_{N_e})}{M_P^3|V'(\varphi_{N_e})|} \simeq \frac{1}{5\sqrt{3}\pi} \times 5.3 \times 10^{-4} \quad (1)$$

is imposed. Here, M_P denotes the reduced Planck scale $2.4 \times 10^{18}\text{GeV}$, V the inflaton potential, V' the derivative of the potential with respect to the inflaton φ , and φ_{N_e} the value of the inflaton field at the time the COBE-scale exited the horizon. Different values of model parameters lead to different reheating processes after those inflations. Successful inflation models must have viable parameter spaces whose resulting reheating processes satisfy all phenomenological constraints such as those from the gravitino problem.

Whether a sufficient e-fold number is obtained for a given model parameter must be examined using the *full* supergravity potential; it must be noted that the inflation sector is not the only one which couples to supergravity. In particular, the hidden sector, which has SUSY-breaking expectation values, is also coupled to supergravity. Then, the inflaton potential receives soft SUSY-breaking contributions, as in the ordinary gravity mediation of the SUSY breaking to the standard-model sector¹. It is true that the gravity-mediated soft SUSY-breaking terms appear with $1/M_P$ -suppression and are tiny. However, the flatness of the slow-roll-inflaton potential requires a tuning at the level of order of the Hubble parameter $H = \sqrt{V/3}/M_P$, which is also a Planck-suppressed quantity. Therefore, the perturbations to the inflaton potential by soft SUSY-breaking terms can be crucial for its flatness. This observation gives new constraints on the viable parameter space of successful inflation.

In this letter, we show how the existence of the soft SUSY-breaking terms constrains the parameter spaces of inflation models. We apply the above-mentioned new generic constraints to a hybrid inflation model in supergravity, as an example, and derive a lower bound on the Yukawa coupling of the inflaton field φ and the inflaton mass m_φ . We show that this new lower bound is highly significant if we take seriously the severest constraint from avoiding the gravitino problems. We also discuss analogous constraints on other inflation models.

¹Ref.[6] pointed out that these SUSY breaking terms can have physical importance.

The slow-roll inflation works well when the slow-roll conditions of the inflaton potential,

$$\epsilon \equiv \left(\frac{V'(\varphi) M_P}{V(\varphi)} \right)^2 \ll 1, \quad (2)$$

$$\eta \equiv \frac{V''(\varphi) M_P^2}{V(\varphi)}, \quad |\eta| \ll 1, \quad (3)$$

are satisfied [7]. Gravity-mediated soft SUSY-breaking terms should not violate the above two conditions. We can robustly say from the condition eq.(3) that

$$|m_{\text{SUSY}}^2| \lesssim |V''| \ll \frac{V}{M_P^2} \simeq H^2, \quad (4)$$

where m_{SUSY} is the soft SUSY-breaking mass. This condition must be satisfied for any inflation that obtains not-so-small e-fold number,² and puts a lower bound on the possible energy scale of inflations for a fixed gravitino mass. We can obtain severer constraints for each model, although they are model dependent. In the following, we take the hybrid inflation model as an example. Constraints on other inflation models are discussed later.

The hybrid inflation models in supergravity are given by the following Kähler potential and superpotential [2]:

$$K = \Phi^\dagger \Phi - \frac{k}{4} \frac{(\Phi^\dagger \Phi)^2}{M_P^2} + \dots \quad (5)$$

$$W = \Phi(\lambda \Psi \bar{\Psi} - \Lambda_I^2). \quad (6)$$

Here, the radial part of the scalar component of the Φ multiplet (φ) plays the roll of the inflaton. Λ_I is the energy scale of the inflation ($V \simeq \Lambda_I^4$). The SUSY contribution to the inflaton potential is given by

$$V(\varphi) = \Lambda_I^4 \left(1 + \frac{\lambda^2}{8\pi^2} \ln\left(\frac{\varphi}{\varphi_c}\right) + \frac{k}{2} \left(\frac{\varphi}{M_P} \right)^2 + \dots \right), \quad (7)$$

where $\varphi_c \equiv \sqrt{2/\lambda} \Lambda_I$ is the value of the inflaton field where the slow-roll inflation ends. The second term comes from the 1-loop renormalization.

Now we expect that the soft SUSY-breaking terms exist,

$$V_{\text{SUSY}}(\varphi) \simeq \frac{1}{2} m_{\text{SUSY}}^2 \varphi^2 + |A| \Lambda_I^2 \varphi, \quad (8)$$

²It is possible that the soft SUSY-breaking terms accidentally cancel the supersymmetric contribution to the inflaton potential and do not ruin the flatness. In this special case, this generic constraint is not necessarily applicable.

where $m_{S\cancel{U}SY}$ is the soft SUSY-breaking mass and A is a dimension-one SUSY-breaking parameter. Combined with the SUSY contribution in eq.(7), effective mass squared of the inflaton is given by $(\Lambda_I^4/M_P^2)k + m_{S\cancel{U}SY}^2$. One thing we notice is that we can change the quadratic term in the total inflation potential by changing the value of k only if

$$\left| \left(\frac{m_{S\cancel{U}SY}^2}{\Lambda_I^4/M_P^2} \right) \right| \lesssim k. \quad (9)$$

Another thing we point out here is that the slow-roll condition of η is always violated unless

$$|m_{S\cancel{U}SY}^2| M_P^2 \ll \Lambda_I^4, \quad (10)$$

which is nothing but eq.(4). The other slow-roll condition of ϵ is also violated unless

$$|A|^2 M_P^2 \ll \Lambda_I^4. \quad (11)$$

If the soft SUSY-breaking mass and the A -parameter are given by the gravitino mass as in the gravity-mediation model, the energy scale of the inflation is bounded from below as

$$(m_{3/2} M_P)^2 \ll \Lambda_I^4. \quad (12)$$

Again, this bound must be satisfied by any inflation which occurred in the thermal history, irrespective of whether it satisfies the COBE-normalization condition eq.(1) or not. That is, this bound is also applicable to inflations that occurred later than the COBE-normalized inflation [8]. The energy scale of any inflation must be higher than the SUSY-breaking scale $\Lambda_{S\cancel{U}SY} \sim \sqrt{m_{3/2} M_P}$.

A few comments on the induced soft SUSY-breaking terms are in order here.

We assume that the soft SUSY-breaking contributions to the inflation potential during the inflation is not so different from those expected from the present vacuum of the hidden sector. Suppose the hidden sector has flat directions. Then it is possible that the field expectation values of the hidden sector fields are not settled at the present value, since eq.(4) holds. In those cases, there is a possibility that the condensation energy of the hidden sector is by far different from the present value, leading to soft SUSY-breaking terms drastically different from those induced by the present vacuum of the hidden sector. However, this possibility is out of the scope of the discussion in the following.

The other limitation of our analysis is that we assume $m_{S\cancel{U}SY} \sim |A| \sim m_{3/2}$ (gravity-mediation model). For the light gravitino mass $m_{3/2} \lesssim 10\text{GeV}$ (gauge-mediation model), the inflaton can acquire soft SUSY-breaking terms whose soft parameters are of order of the

weak scale ($\gg m_{3/2}$), if the inflatons are charged under the standard-model gauge group. For heavy gravitino mass $10\text{TeV} \lesssim m_{3/2}$ (anomaly-mediation model), the inflaton obtains only loop-suppressed soft SUSY-breaking terms, if the inflaton sector is, along with the observed sector, suitably separated from the hidden sector in the Kähler potential. We also neglect these cases, although necessary modification for these two models of SUSY-breaking mediation is straightforward.

Now, let us see explicitly how the generic constraints discussed above put limits on the parameters of the hybrid inflation model. We apply the constraint eq.(9)³ to the inflation that satisfies the COBE-normalization condition eq.(1). The field value φ_{N_e} in eq.(1) is determined by requiring the obtained e-fold number to be suitable:

$$67 + \frac{1}{3} \ln \left(\frac{HT_R}{M_P^2} \right) = N_e = \frac{1}{M_P^2} \int_{\varphi_c}^{\varphi_{N_e}} d\varphi \frac{V(\varphi)}{V'(\varphi)}, \quad (13)$$

where H is the Hubble parameter during the inflation and T_R the reheating temperature after the inflation. The hybrid inflation model given by eqs.(5) and (6) has three parameters, namely Λ_I, k and λ , among which Λ_I can be expressed as a function of k and λ due to the COBE-normalization condition eq.(1).

It was shown in Ref.[9] that the suitable e-fold number is obtained for the parameter space below the thick black solid line described in Fig.1,⁴ which was derived in the absence of the soft SUSY-breaking terms. Furthermore, they showed numerically that for smaller λ and smaller k , the energy scale of the inflation Λ_I becomes lower.

We present an analytical argument and visualize what happens when the Yukawa coupling λ becomes smaller. For sufficiently small $\lambda \ll 3 \times 10^{-3}$ with negligible k ,⁵ or precisely speaking,

$$2.2 \times 10^1 \lambda^{\frac{4}{3}} = 2.2 \times 10^{-7} \left(\frac{\lambda}{10^{-6}} \right)^{\frac{4}{3}} \gg k, \quad (14)$$

Λ_I and φ_{N_e} are solved in terms of λ and k through eq.(1) and eq.(13) as:

$$\left(\frac{\Lambda_I}{M_P} \right) \simeq 1.7 \times 10^{-2} \lambda^{\frac{5}{6}}, \quad (15)$$

³Eq.(12) is satisfied if we use $k \lesssim 1$ and eq.(9).

⁴Assumption on the reheating process does not change this result very much.

⁵This parameter region corresponds to the situation where the following two conditions are satisfied: (1) $\varphi_c \sim \varphi_{N_e} \ll M_P$ holds, and (2) $V'(\varphi) \simeq \Lambda_I^4 ((\lambda^2/8\pi^2)(1/\varphi) + (k\varphi/M_P^2))$ is dominated by the previous term.

and⁶

$$\left(\frac{\varphi_c}{M_P}\right)^2 \simeq 5.8 \times 10^{-4} \lambda^{\frac{2}{3}}, \quad (16)$$

$$\left(\frac{\varphi_{N_e} - \varphi_c}{M_P}\right) \simeq 0.53(61 + \frac{3}{4} \times \ln \lambda) \lambda^{\frac{5}{3}}. \quad (17)$$

Notice that the energy scale of the inflation is determined only by λ because we consider negligibly small k [see eq.(14)]. The exponents of λ in eqs.(15)–(17) can be easily understood by observing that the following three relations must be satisfied:

$$\left(\frac{\varphi_c}{M_P} \propto \lambda^{\frac{1}{3}}\right) = \sqrt{\frac{2}{\lambda}} \left(\frac{\Lambda_I}{M_P}\right) \propto \lambda^{-\frac{1}{2}} \lambda^{\frac{5}{6}} \propto \lambda^{\frac{1}{3}}, \quad (18)$$

$$(N \sim 60 + \text{factor} \times \ln \lambda) \propto \frac{1}{\lambda^2} \left(\frac{\varphi_{N_e} - \varphi_c}{M_P}\right) \left(\frac{\varphi_c}{M_P}\right) \propto \lambda^{-2} \lambda^{\frac{5}{3}} \lambda^{\frac{1}{3}} \propto \lambda^0, \quad (19)$$

$$\left(\frac{\delta\rho}{\rho} \sim 10^{-5}\right) \propto \left(\frac{\Lambda_I}{M_P}\right)^2 \frac{1}{\lambda^2} \left(\frac{\varphi_c}{M_P}\right) \propto \lambda^{\frac{5}{3}} \lambda^{-2} \lambda^{\frac{1}{3}} \propto \lambda^0. \quad (20)$$

After all, we see that Λ_I decreases as $\Lambda_I \propto \lambda^{5/6}$ when λ decreases.

Remember that Λ_I is bounded from below [see eq.(12)] because of the existence of the soft SUSY-breaking mass terms. This means that there exists a lower bound on the coupling λ . Indeed, eq.(9), which says that the parameter k is bounded from below, is now expressed as (for $\lambda \ll 3 \times 10^{-3}$ and eq.(14))

$$k \gtrsim \left(\frac{M_P}{\Lambda_I}\right)^4 \left(\frac{m_{3/2}}{M_P}\right)^2 \simeq 1.2 \times 10^{-5} \left(\frac{m_{3/2}}{240\text{GeV}}\right)^2 \left(\frac{\lambda}{10^{-6}}\right)^{-\frac{10}{3}}. \quad (21)$$

This constraint eq.(21) is described in Fig.1 as the red (dashed) line. The region between the red (dashed) line and the thick black solid line in Fig.1 is the viable parameter space of the hybrid inflation model with a suitable e-fold number, after we take the SUSY-breaking terms into account. It can be seen from the figure, as we stated before, that the possible value of λ has a lower bound:

$$\lambda_{\min.} \simeq 2.0 \times 10^{-6} \left(\frac{m_{3/2}}{240\text{GeV}}\right)^{\frac{3}{7}} \lesssim \lambda, \quad (22)$$

where we have used an empirical relation read off from the figure

$$4.9 \times 10^{-7} \left(\frac{\lambda}{10^{-6}}\right)^{\frac{4}{3}} \gtrsim k \quad (\text{for } 10^{-6} \lesssim \lambda \lesssim 10^{-4}). \quad (23)$$

⁶In eq.(17), we calculate the reheating temperature using eq.(25) as the decay rate of the inflaton. This assumption does not make any essential difference in the following analysis. See also footnote 4.

The energy scale of the inflation at the extreme case of $\lambda \simeq \lambda_{\min.}$ in eq.(22) is

$$\Lambda_I \simeq 3.0 \times 10^{-7} \left(\frac{m_{3/2}}{240\text{GeV}} \right)^{\frac{5}{14}} M_P = 7.2 \times 10^{11} \left(\frac{m_{3/2}}{240\text{GeV}} \right)^{\frac{5}{14}} \text{GeV}, \quad (24)$$

which is slightly larger than the energy scale of the SUSY breaking, $\Lambda_{\text{SUSY}} \equiv 3^{1/4} \sqrt{m_{3/2} M_P} = 3.2 \times 10^{10} \times (m_{3/2}/240\text{GeV})^{1/2} \text{GeV}$, and hence we see that the soft SUSY-breaking terms become important for the smallest value of λ obtained in eq.(22).

Now let us briefly discuss the implication of the lower bound on the parameter λ . The lower bound eq.(22) is far below the *natural* value (order one), or far below the upper bound for a sufficient e-fold number $\lambda \lesssim (1-2) \times 10^{-1}$ [9]. Therefore, the bound seems to have almost no physical importance at first sight. However, a sufficient e-fold number is not the only requirement to the inflation model; reheating after the inflation must lead to a suitable initial condition of the Big-Bang cosmology, and this requirement gives much stronger constraint on λ .

For instance, the requirement from the problem of thermally-produced gravitinos puts an upper bound on the parameter λ [9], under an assumption that the decay rate of the inflaton is given by

$$\Gamma = \frac{N}{8\pi} \frac{m_\phi^3 \left(\frac{\Lambda}{\sqrt{\lambda}} \right)^2}{M_P^4}, \quad (25)$$

where N is the number of the decay modes of the inflaton (of order of 100), and $m_\phi \equiv \sqrt{2\lambda}\Lambda_I$ the inflaton mass. The inflaton mass m_ϕ is proportional to $\lambda^{4/3}$, and the reheating temperature is proportional to $\lambda^{7/3}$. It is shown in Ref.[9] that $\lambda \lesssim 10^{-3}$ is necessary for $T_R \lesssim 10^6 \text{GeV}$. This upper bound on T_R is derived to avoid the thermal overproduction of gravitinos [10].

Recently, it has been discussed that the gravitinos can also be produced non-thermally during the matter-dominated era before the reheating, and that the abundance of these gravitinos can be larger than that of the gravitinos produced thermally after the reheating. There are two different estimations of the yield (*i.e.* the ratio of the number density $n_{3/2}$ to the entropy s) of the non-thermally produced gravitinos [11, 12]. Among these two, the estimation of Ref.[12] predicts a larger value of the yield, and hence leads to a stronger

constraint on the successful inflation models. The estimation is given by⁷

$$\frac{n_{3/2}}{s} \sim 10^{-2} \frac{m_\varphi^3 T_R}{(m_{3/2} M_P)^2}. \quad (26)$$

Under the assumption of the inflaton decay rate eq.(25), eq.(26) leads to

$$\frac{n_{3/2}}{s} \sim 3.1 \times 10^{-12} \left(\frac{M_P}{m_{3/2}} \right)^2 \lambda^{\frac{19}{3}} \quad (27)$$

for $\lambda \ll 3 \times 10^{-3}$ and negligibly small k (*i.e.* eq.(14)). The cosmological upper bound on the yield of the gravitino ($n_{3/2}/s \lesssim 10^{-12}$ [10]) leads to the following upper bound on λ ;

$$\lambda \lesssim 7.4 \times 10^{-6} \left(\frac{m_{3/2}}{240 \text{ GeV}} \right)^{\frac{6}{19}}. \quad (28)$$

Therefore, even if we adopt the severest estimation of the non-thermal production of the gravitinos, there exists a narrow (but not vanishing) parameter space between the upper bound from avoiding the gravitino overproduction in eq.(28) and the lower bound from the soft SUSY-breaking terms eq.(22). In conclusion, the hybrid inflation model has a viable parameter space where sufficient e-fold number is obtained and at the same time where the yield of the gravitinos produced thermally and non-thermally is sufficiently suppressed.

We have pointed out that the existence of the soft SUSY-breaking terms from the hidden sector can violate the required flatness of the inflation potential, and hence can restrict the viable parameter space of the model. Our consideration has led to the existence of a lower bound on the parameter λ in a hybrid inflation model. Analyses for other inflation models are straightforward, and we do not repeat them here; we only describe the results. The consequences of the gravitino problem can also be derived easily.

For chaotic inflation model [1], our discussion does not give any new information, because the COBE-normalization condition sets the inflaton mass to be of order of 10^{13} GeV , which is far above the soft SUSY-breaking mass.

For the new inflation model in the 2nd reference of [3],⁸ we need to consider two cases, according to which term dominates in the first derivative of the inflation potential $V'(\varphi_{N_e})$,

⁷While preparing this paper for publication, we received the preprint [13] which says that the gravitinos produced while the Hubble parameter is larger than the gravitino mass remain as inflatinos under the time evolution and do not turn into the present day gravitinos in a model they considered. In the case this statement also holds in a hybrid inflation model, the yield of the non-thermally produced gravitinos reduces significantly from eq.(26).

⁸See Ref. [3] for the notations.

i.e., the $\varphi_{N_e}^1$ term or the $\varphi_{N_e}^{n-1}$ term. In the $\varphi_{N_e}^1$ -dominant parameter region ($1/[(n-2)N_e + (n-1)] < k$), the energy scale of the inflation Λ_I is written in terms of k and g as

$$\Lambda_I \simeq M_P \cdot (\sqrt{2}c)^{\frac{n-2}{2(n-3)}} \cdot k^{\frac{n-1}{2(n-3)}} \cdot (ng)^{-\frac{1}{2(n-3)}} \cdot e^{-\frac{n-2}{2(n-3)}kN}, \quad (29)$$

where $c \equiv 5.3 \times 10^{-4}$ in eq.(1). Our constraint eq.(4) gives a lower bound on $k\Lambda_I^4(k, g)$ for a fixed gravitino mass and fixed k , and as a consequence, g is bounded from above for fixed k (since $\Lambda_I \propto g^{-\frac{1}{2(n-3)}}$). Then, the inflaton mass $m_\varphi \propto g^{-\frac{2}{n(n-3)}}$ is bounded from below for fixed k , and this lower bound takes its minimum value at $k \simeq 0.1$;⁹ it is of order of 10^7 GeV for $n = 4$. In the $\varphi_{N_e}^{n-1}$ -dominant parameter region ($0 < k < 1/[(n-2)N_e + (n-1)]$), the inflaton mass is always larger than the smallest inflaton mass in the $\varphi_{N_e}^1$ -dominant parameter region.

For the new inflation model in the 1st reference of [3], the potential is the same as that discussed in the previous paragraph, with one extra constraint between Λ_I and g imposed. As a result, we can write g (and hence the inflaton mass $m_\varphi \simeq ng^{1/n}\Lambda_I^{2(n-1)/n}/M_P^{(n-2)/n}$) in terms of k and n (in the $\varphi_{N_e}^1$ -dominant region), or n (in the $\varphi_{N_e}^{n-1}$ -dominant region). We see that the lower bound on the inflaton mass is of order of 10^7 GeV , again at the observed upper bound on k , although this bound has nothing to do with the existence of the soft SUSY-breaking terms in this model.

For the topological inflation model [4], the energy scale of inflation Λ_I can be written solely in terms of $\kappa \equiv 2g + k - 1$, where g is the Yukawa coupling in the superpotential and k is the four-point coupling in the Kähler potential; $\Lambda_I = (c\kappa)^{1/2}M_P e^{-\kappa N/2}$, where $c \equiv 5.3 \times 10^{-4}$. Our constraint eq.(4) gives a lower bound on $\kappa\Lambda_I^4(\kappa)$ for a fixed gravitino mass, and hence a bound on κ , $10^{-8} \lesssim \kappa \lesssim 0.5$. However, eq.(4) does not give any constraint on g , so thermal and non-thermal production of gravitinos can be sufficiently suppressed by taking g small enough.¹⁰

For a D-term inflation model [5], the potential of this model is nothing but that of the hybrid inflation model eq.(7) with $k = 0$, λ replaced by the gauge coupling g , and Λ_I^4 by $\frac{1}{2}(g^2\xi^4)$, where ξ^2 is the Fayet-Iliopoulos parameter. The same analysis developed in this paper is applicable, and the gauge coupling g is bounded from below. However, the inflaton mass still depends upon the Yukawa coupling λ in the superpotential, which is not constrained by the presence of the hidden sector. Therefore, no new results on the reheating process are derived.

⁹The observational constraint on the spectral index $n_s \gtrsim 0.8$ puts an upper bound on k (since $n_s = 1 - 2k$).

¹⁰Inflaton mass m_φ is written as $m_\varphi = 2\sqrt{g}\Lambda_I^2/M_P$ in this model.

Acknowledgments

We thank Tsutomu Yanagida for his careful reading of this manuscript, precious comments, and continuous encouragements. T.W. thanks the Japan Society for the Promotion of Science for financial support.

References

- [1] H. Murayama, H. Suzuki, T. Yanagida and J. Yokoyama, Phys.Rev. **D50**, 2356-2360 (1994);
M. Kawasaki, Masahide Yamaguchi and T. Yanagida, Phys.Rev.Lett. **85**, 3572-3575 (2000).
- [2] E. J. Copeland, A. R. Liddle, D. H. Lyth, E. D. Stewart and D. Wands, Phys.Rev. **D49**, 6410 (1994);
A.D.Linde, Phys.Lett. **B259**, 38 (1991);
G. Dvali, R. Schaefer and Q. Shafi, Phys.Rev.Lett. **73**, 1886 (1994);
G. Lazarides, R. Schaefer and Q. Shafi, Phys.Rev. **D56**, 1324 (1997);
A. Linde and A. Riotto, Phys.Rev. **D56**, 1841 (1997).
- [3] K.-I. Izawa and T. Yanagida, Phys.Lett. **B393**, 331-336 (1997);
Ref.[9].
- [4] K.I. Izawa, M. Kawasaki and T. Yanagida, Prog.Theor.Phys. **101**, 1129-1133 (1999);
M. Kawasaki, N. Sakai, Masahide Yamaguchi and T. Yanagida, Phys.Rev. **D62**, 123507 (2000).
- [5] E. Stewart, Phys.Rev. **D51**, 6847 (1995);
E. Halyo, Phys.Lett. **B387**, 43 (1996);
P. Binétruy and G. Dvali, Phys.Lett. **B388**, 241 (1996).
- [6] W. Buchmüller, L. Covi and D. Delépine, Phys.Lett. **B491**, 183-189 (2000).
- [7] E. W. Kolb and M. S. Turner, “The Early Universe,” *Redwood City, USA: Addison-Wesley (1990) 547 p. (Frontiers in physics, 69)*.
- [8] M. Kawasaki, N. Sugiyama and T. Yanagida, Phys.Rev. **D57**, 6050-6056 (1998);
T. Kanazawa, M. Kawasaki, N. Sugiyama and T. Yanagida, Phys.Rev. **D61**, 023517 (2000);
T. Kanazawa, M. Kawasaki, N. Sugiyama and T. Yanagida, astro-ph/0006445.
- [9] T. Asaka, K. Hamaguchi, M. Kawasaki and T. Yanagida, Phys.Rev. **D61**, 083512 (2000).
- [10] E. Holtmann, M. Kawasaki, K. Kohri, T. Moroi, Phys.Rev. **D60**, 023506 (1999).
- [11] G. Giudice, A. Riotto and I. Tkachev, JHEP **9908**, 009 (1999), JHEP **9911**, 036 (1999).
- [12] D. H. Lyth, Phys.Lett. **B488**, 417-422 (2000), Phys.Lett. **B476**, 356-362 (2000);
H. B. Kim and D. H. Lyth, hep-ph/0011262.

[13] H.P. Nilles, M. Peloso and L. Sorbo, hep-ph/0102264.

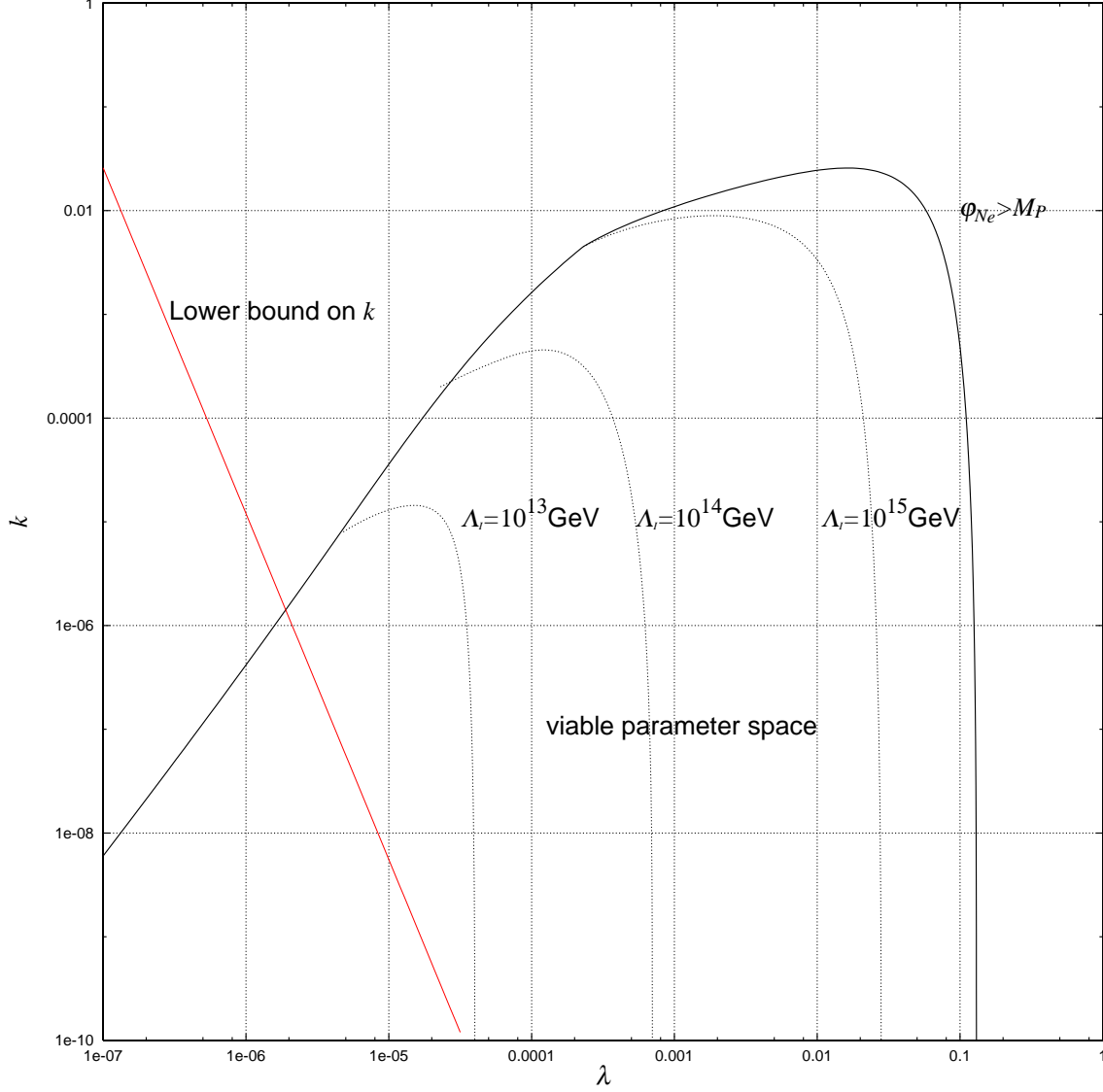


Figure 1: The viable parameter space of the COBE-normalized hybrid inflation model. The red (dashed) line is the lower bound on k from the SUSY-breaking soft terms. The thick black solid line is the upper bound on k from the requirement $\varphi_{N_e} < M_P$. We can see from this that λ is lower bounded, $\lambda \gtrsim 2.0 \times 10^{-6}$. We also show the contours for fixed Λ_I . The lines except the red (dashed) line are cited (with some modification) from Ref.[9]. We thank the authors for permission.